Dark Energy and Dark Matter from Quantum Gravity

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In this talk

☞ I emphasize that noncommutative (NC) spacetime necessarily implies emergent spacetime if spacetime at microscopic scales should be viewed as NC.

☞ The emergent gravity from NC $U(1)$ gauge theory is the large N duality and the emergent spacetime picture admits a background-independent formulation of quantum gravity.

☞ In order to understand NC spacetime correctly, we need to deactivate the thought patterns that we have installed in our brains and taken for granted for so many years.

☞ Emergent gravity predicts the existence of the dark composition of our Universe.

NC spacetime introduces the gauge-gravity duality

1. Recall that quantum mechanics is mechanics on NC phase space whose coordinate generators satisfy the commutation relation

$$[x^i, p_j] = i\hbar\delta^i_j$$

2. The mathematical structure of NC spacetime is essentially the same as the NC phase space in quantum mechanics:

$$[y^\mu, y^\nu] = i\theta^{\mu\nu},$$

where $p_\mu = B_{\mu\nu}\, y^\nu$ and $B_{\mu\nu} \equiv (\theta^{-1})_{\mu\nu}$.

3. Everything on NC spacetime bears some analogy with quantum mechanics:

NC phase space $\Longrightarrow$ Wave-particle duality

NC spacetime $\Longrightarrow$ Gauge-gravity duality

NC spacetime necessarily implies emergent spacetime

1. Recall that $f(x+a) = U(a)^\dagger f(x) U(a)$ where $U(a) = e^{-ip\cdot a/\hbar}$,
   and so every point on NC space is unitarily equivalent.
   There is no space but an algebra $\mathcal{A}_\theta$ only. Thus the NC *space* is a misnomer.

2. NC space introduces a separable Hilbert space $\mathcal{H}$ and so $\mathcal{H}$ has a countable basis.
   Dynamical variables become operators acting on the Hilbert space.

3. NC spacetime implies a paradigm shift: Geometry $\longrightarrow$ Algebra
   Hilbert space $\mathcal{H}$: dynamical variables $\longrightarrow$ $N \times N$ matrices
   where $N = \dim(\mathcal{H}) \rightarrow \infty$.

4. Large N duality or gauge/gravity duality such as the AdS/CFT correspondence
   is an inevitable consequence of the NC spacetime.

5. NC spacetime admits a (dynamical) diffeomorphism symmetry which precisely
   acts as the novel form of the equivalence principle for electromagnetic force.

NC Fields As Large N Matrices

Consider a two-dimensional NC space

$$[x, y] = i\theta \quad \Leftrightarrow \quad [a, a^{\dagger}] = 1 \text{ where } a = \frac{x+iy}{\sqrt{2\theta}}.$$

Since $\mathcal{H} = \{\, |n\rangle\, ;\ n = 0,1,\cdots,\infty\}$ and $\sum_{n=0}^{\infty} |n\rangle\langle n| = 1_{\mathcal{H}}$, for $\phi_1,\ \phi_2 \in \mathcal{A}_\theta$,

$$\phi_1(x,y) = \sum_{n,m=0}^{\infty} |n\rangle\langle n|\phi_1(x,y)\,|m\rangle\langle m| \equiv M_{nm}|n\rangle\langle m|,$$

$$\phi_2(x,y) = \sum_{n,m=0}^{\infty} |n\rangle\langle n|\phi_2(x,y)\,|m\rangle\langle m| \equiv N_{nm}|n\rangle\langle m|,$$

$$(\phi_1 \star \phi_2)(x,y) = \sum_{n,l,m=0}^{\infty} |n\rangle\langle n|\phi_1(x,y)|l\rangle\langle l|\phi_2(x,y)\,|m\rangle\langle m| = M_{nl}N_{lm}|n\rangle\langle m|,$$

NC fields $\phi_a(x,y)$ in $\mathcal{A}_\theta$= adjoint operators acting on a separable Hilbert space $\mathcal{H} = N \times N$ matrices in $End(\mathcal{H}) \equiv \mathcal{A}_N$ with $N = \dim(\mathcal{H}) \to \infty$.

Ordering in $\mathcal{A}_\theta$= ordering in $\mathcal{A}_N$ and $Tr_N = Tr_{\mathcal{H}} = \int \frac{dxdy}{2\pi\theta}$.

Large N gauge theory from NC U(1) gauge theory

Consider a (d+2n)-dimensional NC $U(1)$ gauge theory on $\mathbb{R}^d_C \times \mathbb{R}^{2n}_{NC}$ whose coordinates are $X^M = (x^\mu, y^a)$, $M = 0,1,\cdots,D-1$, $\mu = 0,1,\cdots,d-1$, $a = 1,\cdots,2n$ where

$$[y^a, y^b] = i\theta^{ab}.$$

The D=(d+2n)-dimensional $U(1)$ connections are split as

$$D_M(X) = \partial_M - i\hat{A}_M(x,y) = \big(D_\mu, D_a\big)(x,y)$$

where $\partial_a \equiv \mathrm{ad}_{p_a} = -i[p_a, \cdot]$ with $p_a = B_{ab}y^b$ and

$$D_a\,(x,y) = -i\big(p_a + \hat{A}_a(x,y)\big) \equiv -i\phi_a(x,y).$$

Using the matrix representation $\mathcal{A}_\theta \to \mathcal{A}_N$ by

$$\Xi(x,y) \mapsto \Xi(x) \in \mathcal{A}_N,$$

the D-dimensional NC $U(1)$ gauge theory is exactly mapped to the d-dimenisonal $U(N)$ Yang-Mills theory

$$\begin{aligned} S &= -\frac{1}{4G_{YM}^2}\int d^D X\,\big(\hat{F}_{MN} - B_{MN}\big)^2 \\ &= -\frac{1}{g_{YM}^2}\int d^d\mathrm{x}\,\mathrm{Tr}\big(\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} + \tfrac{1}{2}D_\mu\phi_a D^\mu\phi_a - \tfrac{1}{4}[\phi_a,\phi_b]^2\big) \end{aligned}$$

where $B_{MN} = \begin{pmatrix} 0 & 0 \\ 0 & B_{ab} \end{pmatrix}$.

d=0: IKKT, d=1: BFSS, d=2: DVV, ⋯ , d=4: AdS/CFT, ⋯ .

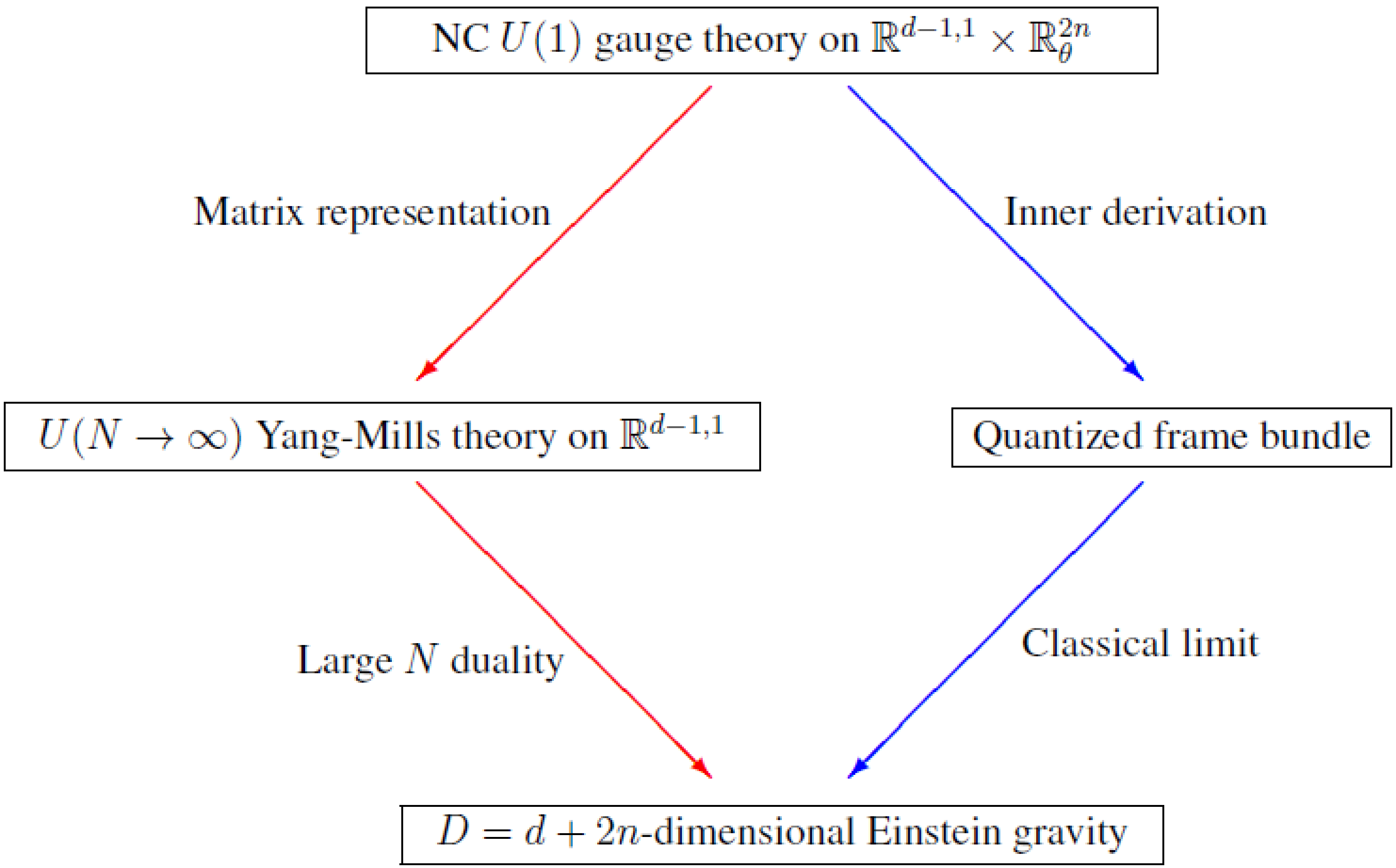

Figure 1: Flowchart for emergent gravity

NC spacetime as NC Coulomb branch

I will make an important observation that NC spacetime arises as a vacuum solution in the Coulomb branch of a large N gauge theory to demonstrate the large N duality.

1. The conventional choice of vacuum in the Coulomb branch of $U(N)$ Yang-Mills theory is given by

$$[\phi_a, \phi_b]|_{\text{vac}} = 0 \qquad \Rightarrow \qquad \langle \phi_a \rangle_{\text{vac}} = \text{diag}\big((\alpha_a)_1, (\alpha_a)_2, \cdots, (\alpha_a)_N\big)$$

   In this case the U(N) gauge symmetry is broken to $U(1)^N$.

2. If we consider the $N \to \infty$ limit, the large N limit opens a new phase of the Coulomb branch given by

$$[\phi_a, \phi_b]|_{\text{vac}} = -iB_{ab} \qquad \Rightarrow \qquad \langle \phi_a \rangle_{\text{vac}} = p_a \equiv B_{ab} y^b$$

   where the vacuum moduli $y^a \in \mathbb{R}^{2n}_{NC}$ satisfy the Moyal-Heisenberg algebra.

3. Suppose that fluctuations around the NC vacuum take the form

$$D_\mu = \partial_\mu - i\widehat{A}_\mu(x, y), \qquad \phi_a = p_a + \widehat{A}_a(x, y)$$

   The adjoint scalar fields now obey the deformed algebra given by

$$[\phi_a, \phi_b] = -i(B_{ab} - \widehat{F}_{ab})$$

Large N duality from NC spacetime

Plugging the fluctuations into the *d*-dimensional $U(N \to \infty)$ Yang-Mills theory, we get the $D = (d + 2n)$-dimensional NC $U(1)$ gauge theory and thus arrive at the reversed version of the equivalence

$$S = -\frac{1}{g_{YM}^2}\int d^d\mathrm{x}\,\mathrm{Tr}\big(\tfrac{1}{4}F_{\mu\nu}F^{\mu\nu} + \tfrac{1}{2}D_\mu\phi_a D^\mu\phi_a - \tfrac{1}{4}[\phi_a,\phi_b]^2\big)$$
$$= -\frac{1}{4G_{YM}^2}\int d^D X\big(\hat{F}_{MN} - B_{MN}\big)^2$$

where $X^M = (x^\mu, y^a)$ are D-dimensional coordinates and D-dimensional connections are defined by

$$D_M(X) = \partial_M - i\hat{A}_M(x,y) = \big(D_\mu, D_a\big)(x,y)$$

whose field strength is given by

$$\hat{F}_{MN}(X) = \partial_M\hat{A}_N - \partial_N\hat{A}_M - i\big[\hat{A}_M,\hat{A}_N\big]_\star.$$

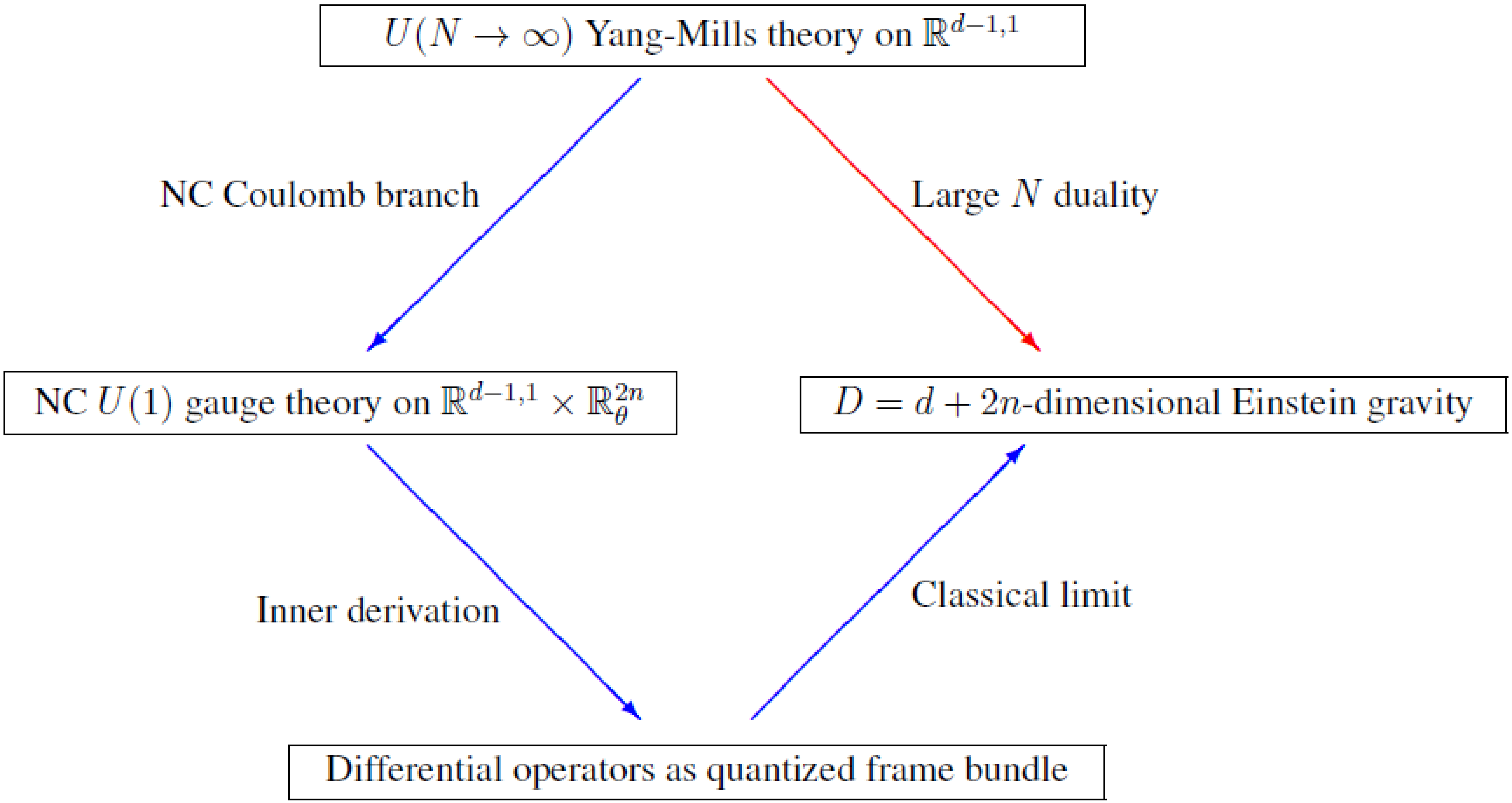

Figure 2: Flowchart for large N duality

Inner/Outer automorphism

Therefore there should be some way to map the NC $U(1)$ gauge theory to the Einstein gravity according to the (conjectural) large N duality. To be more specific, consider the inverse metric in Einstein gravity given by

$$\left(\frac{\partial}{\partial s}\right)^2 = E_A \otimes E_A = g^{MN}(X)\partial_M \otimes \partial_N$$

where $E_A = E_A^M(X)\partial_M$ are orthonormal frames on the tangent bundle $T\mathcal{M}$ of a D-dimensional spacetime manifold $\mathcal{M}$. In order to complete the gauge-gravity duality in Figs. 1 and 2, it is thus necessary to realize the vector fields $E_A = E_A^M(X)\partial_M \in \Gamma(T\mathcal{M})$ in terms of NC $U(1)$ gauge fields.

A decisive clue is coming from the fact that the NC $\star$-algebra $\mathcal{A}_\theta$ generated by the Moyal-Heisenberg algebra always admits a nontrivial inner automorphism $\mathfrak{I}$. The infinitesimal generators of $\mathfrak{I}$ form an inner derivation defined by the adjoint operation

$$\mathcal{A}_\theta^d \to \mathfrak{D}^d : f \mapsto \mathrm{ad}_f = -i[f,\ \cdot]$$

for any $f \in \mathcal{A}_\theta^d \equiv \mathcal{A}_\theta\big(C^\infty(\mathbb{R}^{d-1,1})\big) \cong \mathcal{A}_\theta \times C^\infty(\mathbb{R}^{d-1,1})$

Definitely the derivation $\mathfrak{D}^d$ is a Lie algebra homomorphism, i.e.,

$$\mathrm{ad}_{[f,g]} = i[\mathrm{ad}_f, \mathrm{ad}_g]$$

Vielbeins from inner derivation

Consider the derivation algebra generated by the dynamical variables defined by

$$\widehat{V}_A = \{i\,\mathrm{ad}_{D_A} = [D_A,\ \cdot\ ] | D_A(x,y) = (D_\mu, D_a)(x,y)\} \in \mathfrak{D}^d$$

where $D_A(x,y) = -i\phi_A(x,y)$. In a large-distance limit, i.e. $|\theta| \to 0$, one can expand the NC vector fields using the explicit form of the Moyal $\star$-product. The result takes the form

$$\widehat{V}_A = V_A^M(x,y)\frac{\partial}{\partial Y^M} + \sum_{p=2}^{\infty} V_A^{a_1\cdots a_p}(x,y)\frac{\partial}{\partial y^{a_1}}\cdots\frac{\partial}{\partial y^{a_p}} \in \mathfrak{D}^d$$

where $V_A^\mu = \delta_A^\mu$. Thus the Taylor expansion of NC vector fields generates an infinite tower of the so-called polyvector fields. Note that the leading term gives rise to the ordinary vector fields that will be identified with a frame basis associated with the tangent bundle $T\mathcal{M}$ of an emergent spacetime manifold $\mathcal{M}$.

Let us truncate the above polyvector fields to ordinary vector fields given by

$$\mathfrak{X}(\mathcal{M}) = \left\{ V_A = V_A^M(x,y)\frac{\partial}{\partial X^M} | A, M = 0, 1, \cdots, D-1 \right\}$$

where $X^M = (x^\mu, y^a)$ are local coordinates on a D-dimensional emergent *Lorentzian* manifold $\mathcal{M}$.

Emergent gravity from NC spacetime

The orthonormal vielbeins on $T\mathcal{M}$ are then defined by the relation

$$V_A = \lambda E_A \in \Gamma(T\mathcal{M})$$

or on $T^*\mathcal{M}$

$$v^A = \lambda^{-1} e^A \in \Gamma(T^*\mathcal{M}).$$

The conformal factor $\lambda \in C^\infty(\mathcal{M})$ is determined by the volume preserving condition

$$\mathcal{L}_{V_A}\nu_t = \big(\nabla \cdot V_A + (d-2n)V_A \ln\lambda\big)\nu_t = 0, \qquad \forall A = 0, 1, \cdots, D-1,$$

where

$$\nu_t \equiv d^d x \wedge \nu = \lambda^2 d^d x \wedge v^1 \wedge \cdots \wedge v^{2n}$$

is a D-dimensional volume form on $\mathcal{M}$.

In the end, the Lorentzian metric on a *D*-dimensional spacetime manifold is given by

$$\begin{aligned} ds^2 &= \mathcal{G}_{MN}(X) dX^M \otimes dX^N = e^A \otimes e^A \\ &= \lambda^2 v^A \otimes v^A = \lambda^2 \big(\eta_{\mu\nu} dx^\mu dx^\nu + v^a_b v^a_c (dy^b - \mathbf{A}^b)(dy^c - \mathbf{A}^c)\big) \end{aligned}$$

where $\mathbf{A}^b := A^b_\mu(x,y)dx^\mu$. Therefore the NC field theory representation of the *d*-dimensional large *N* gauge theory in the NC Coulomb branch provides a powerful machinery to identify gravitational variables dual to large *N* matrices.

NC Electromagnetism from Matrix Model

Let us start with a zero-dimensional matrix model with a bunch of $N \times N$ Hermitian matrices, $\{\phi_a \in \mathcal{A}_N | a = 1, \cdots, 2n\}$, whose action is given by

$$S = -\frac{1}{4} \sum_{a,b=1}^{2n} \mathrm{Tr}\, [\phi_a, \phi_b]^2.$$

We require that the matrix algebra $\mathcal{A}_N$ is associative, which leads to the Jacobi identity

$$[\phi_a, [\phi_b, \phi_c]] + [\phi_b, [\phi_c, \phi_a]] + [\phi_c, [\phi_a, \phi_b]] = 0.$$

We also assume the action principle, from which we yield the equation of motion

$$\sum_{b=1}^{2n} [\phi_b, [\phi_a, \phi_b]] = 0.$$

First suppose that the vacuum configuration of $\mathcal{A}_N$ is given by

$$\langle\phi_a\rangle_{\text{vac}} = p_a \in \mathcal{A}_N,$$

An obvious solution in the limit $N \to \infty$ is given by the Moyal-Heisenberg algebra

$$[p_a, p_b] = -iB_{ab},$$

where $(B_{ab}) = -L_P^{-2}(\mathbf{1}_n \otimes i\sigma^2)$ is a $2n \times 2n$ constant symplectic matrix.
A general solution is generated by considering all possible deformations of the Moyal-Heisenberg algebra. They take the form

$$\phi_a = p_a + \widehat{A}_a \in \mathcal{A}_N,$$

obeying the deformed algebra

$$[\phi_a, \phi_b] = -i(B_{ab} - \widehat{F}_{ab}),$$

where $\widehat{F}_{ab} = \partial_a\widehat{A}_b - \partial_b\widehat{A}_a - i[\widehat{A}_a, \widehat{A}_b] \in \mathcal{A}_N$ with the definition $\partial_a \equiv \mathrm{ad}_{p_a} = -i[p_a, \cdot]$.

For the general matrix $\phi_a \in \mathcal{A}_N$ to be a solution the set of matrices $\widehat{F}_{ab} \in \mathcal{A}_N$, must obey the following equations

$$\widehat{D}_a \widehat{F}_{bc} + \widehat{D}_b \widehat{F}_{ca} + \widehat{D}_c \widehat{F}_{ab} = 0,$$

$$\sum_{b=1}^{2n} \widehat{D}_b \widehat{F}_{ab} = 0,$$

where

$$\widehat{D}_a \widehat{F}_{bc} \equiv \mathrm{ad}_{\phi_a} \widehat{F}_{bc} = -i[\phi_a, \widehat{F}_{bc}] = -[\phi_a, [\phi_b, \phi_c]].$$

Let us apply the Lie algebra homomorphism to yield

$$X_{D_A F_{BC}} = [V_A, [V_B, V_C]] \in \Gamma(TM)$$

It is then straightforward to get the following correspondence

$$D^B F_{AB} = 0 \quad \Leftrightarrow \quad [V^B, [V_A, V_B]] = 0,$$

$$D_A F_{BC} + \text{cyclic} = 0 \quad \Leftrightarrow \quad [V_A, [V_B, V_C]] + \text{cyclic} = 0.$$

Note that the torsion T and curvature R are multi-linear differential operators

$$\begin{aligned} T(X,Y) &= \nabla_X Y - \nabla_Y X - [X,Y], \\ R(X,Y)Z &= [\nabla_X, \nabla_Y]Z - \nabla_{[X,Y]}Z, \end{aligned}$$

where X, Y and Z are vector fields on M. It is easy to derive the relation

$$\begin{aligned} T(V_A, V_B) &= \lambda^2 T(E_A, E_B), \\ R(V_A, V_B)V_C &= \lambda^3 R(E_A, E_B)E_C. \end{aligned}$$

After imposing the torsion free condition $T(E_A, E_B) = 0$,

it is easy to derive the identity below

$$R(E_A, E_B)E_C + \text{cyclic} = \lambda^{-3}\Big([V_A, [V_B, V_C]] + \text{cyclic}\Big).$$

Therefore we see that the Bianchi identity for NC $U(1)$ gauge fields is equivalent to the first Bianchi identity for the Riemann curvature tensors, i.e.,

$$D_A F_{BC} + \text{cyclic} = 0 \qquad \Leftrightarrow \qquad R(E_A, E_B)E_C + \text{cyclic} = 0.$$

The mission for the equations of motion is more involved.
So let us focus on four dimensions, i.e., $n = 2$.
Basically we are expecting the Einstein equations

$$D^B F_{AB} = 0 \qquad \Leftrightarrow \qquad R_{AB} = 8\pi G\Big(T_{AB} - \frac{1}{2}\delta_{AB}T\Big).$$

After a straightforward but tedious calculation, we get a remarkably simple but cryptic result

$$R_{AB} = -\frac{1}{\lambda^2}\Big[g_D^{(+)a} g_D^{(-)\dot{b}}\Big(\eta^a_{AC}\bar{\eta}^{\dot{b}}_{BC} + \eta^a_{BC}\bar{\eta}^{\dot{b}}_{AC}\Big) - g_C^{(+)a} g_D^{(-)\dot{b}}\Big(\eta^a_{AC}\bar{\eta}^{\dot{b}}_{BD} + \eta^a_{BC}\bar{\eta}^{\dot{b}}_{AD}\Big)\Big].$$

To get the above result, we have introduced the structure equation of vector fields $V_A \in \Gamma(TM)$

$$[V_A, V_B] = -g_{AB}{}^C V_C$$

and the canonical decomposition

$$g_{ABC} = g_C^{(+)a}\eta^a_{AB} + g_C^{(-)\dot{a}}\bar{\eta}^{\dot{a}}_{AB}.$$

First it is convenient to decompose the energy-momentum tensor into two parts

$$
\begin{aligned}
8\pi G T^{(M)}_{AB} &= -\frac{1}{\lambda^2}\Big(g_{ACD}g_{BCD} - \frac{1}{4}\delta_{AB}g_{CDE}g_{CDE}\Big), \\
8\pi G T^{(L)}_{AB} &= \frac{1}{2\lambda^2}\Big(\rho_A\rho_B - \Psi_A\Psi_B - \frac{1}{2}\delta_{AB}(\rho_C^2 - \Psi_C^2)\Big),
\end{aligned}
$$

where $\rho_A \equiv g_{BAB}$ and $\Psi_A \equiv -\frac{1}{2}\varepsilon^{ABCD}g_{BCD}$.

A close inspection reveals that the first one is the Maxwell energy-momentum tensor given by

$$
T^{(em)}_{ab} = \frac{\hbar^2 c^2}{g^2_{YM}}\Big(F_{ac}F_{bc} - \frac{1}{4}\delta_{ab}F_{cd}F_{cd}\Big),
$$

but the second one seems to be very mysterious.

In order to descry closer aspects of the second energy-momentum tensor, let us consider the following decomposition

$$
\begin{aligned}
\rho_a\rho_b &= \frac{1}{4}\delta_{ab}\rho_c^2 + \Big(\rho_a\rho_b - \frac{1}{4}\delta_{ab}\rho_c^2\Big), \\
\Psi_a\Psi_b &= \frac{1}{4}\delta_{ab}\Psi_c^2 + \Big(\Psi_a\Psi_b - \frac{1}{4}\delta_{ab}\Psi_c^2\Big).
\end{aligned}
$$

In a long wavelength limit, the quadruple modes can be ignored and then it behaves like a cosmological constant

$$T_{ab}^{(L)} = -\frac{c^4 R}{32\pi G}\delta_{ab}$$

where $R = \frac{1}{2\lambda^2}(\rho_a\rho_b + \Psi_a\Psi_b)\delta^{ab}$. So it may be related to dark energy/dark matter.

In order to get a corresponding result in (3+1)-dimenisonal Lorentzian spacetime, let us take the analytic continuation defined by $x^4 = ix^0$
Under this Wick rotation,

$$\delta_{AB} \to \eta_{AB}, \qquad \Psi_A \to i\Psi_A$$

the so-called Poisson-Liouville energy-momentum (PLEM) tensor is given by

$$T_{\mu\nu}^{(L)} = \frac{1}{16\pi G_4 \lambda^2}\Big(\rho_\mu\rho_\nu + \Psi_\mu\Psi_\nu - \frac{1}{2}g_{\mu\nu}(\rho_\lambda^2 + \Psi_\lambda^2)\Big),$$

Note that ρ_μ and Ψ_μ are four vectors and random fluctuations in nature.
So they are classified into two classes:

(ρ_μ, Ψ_μ): spacelike vectors, i.e., $\rho_\mu\rho_\nu\, g^{\mu\nu} > 0$, etc.
(ρ_μ, Ψ_μ): timelike vectors, i.e., $\rho_\mu\rho_\nu\, g^{\mu\nu} < 0$, etc.

Given a timelike unit vector field u^μ, *i.e.*, $u^\mu u_\mu = -1$, the Raychaudhuri equation in four dimensions is given by

$$\dot{\Theta} - \dot{u}^\mu_{;\mu} + \Sigma_{\mu\nu}\Sigma^{\mu\nu} - \Omega_{\mu\nu}\Omega^{\mu\nu} + \frac{1}{3}\Theta^2 = -R_{\mu\nu}u^\mu u^\nu$$

where

$$-R_{\mu\nu}u^\mu u^\nu = -\frac{1}{2\lambda^2}u^\mu u^\nu(\rho_\mu\rho_\nu + \Psi_\mu\Psi_\nu)$$

Suppose that all the terms except the expansion evolution, $\dot{\Theta}$, in the Raychaudhuri equation vanish or become negligible. In this case the Raychaudhuri equation reduces to

$$\dot{\Theta} = -\frac{1}{2\lambda^2}u^\mu u^\nu(\rho_\mu\rho_\nu + \Psi_\mu\Psi_\nu).$$

Note that $\dot{\Theta} \approx \frac{R}{4}$ and $R < 0$ when ρ_μ and Ψ_μ are timelike while $R > 0$ when ρ_μ and Ψ_μ are spacelike.

$\dot{\Theta} \approx \frac{R}{4} < 0$ for timelike fluctuations

$\dot{\Theta} \approx \frac{R}{4} > 0$ for spacelike fluctuations

By a simple dimensional argument, it is natural to assess that $R \sim \frac{1}{L_H^2}$
Then the PLEM tensor is given by

$$T_{\mu\nu} \approx -\frac{1}{L_P^2 L_H^2} \approx (10^{-3} eV)^4$$

if we identify L_H with the size of cosmic horizon of our observable universe,

$L_H \sim 1.3 \times 10^{26}$ m,

This extended (nonlocal) energy is in good agreement with the observed value of current dark energy.

I quote the footnote in one of my papers.

[7] In the Lorentzian signature, the sign of the Ricci scalar R depends on whether fluctuations are spacelike ($R > 0$) or timelike ($R < 0$) [12, 5]. In consequence the spacelike perturbations act as a repulsive force whereas the timelike ones act as an attractive force. When considering the fact that the fluctuations in (79) are random in nature and we are living in (3+1) (macroscopic) dimensions, the ratio of the repulsive and attractive components will end in $\frac{3}{4} : \frac{1}{4} = 75 : 25$. Is it outrageous to conceive that this ratio curiously coincides with the dark composition of our universe ?